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On the indistinguishibility of Majorana- from Dirac-neutrino propagation in a stellar medium

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Abstract

We study, in the framework of the Standard Model, the propagation of (pure) Majorana neutrinos in a typical stellar medium and show that Majorana neutrino matter oscillations are completely indistinguishable from Dirac ones, even if in the case of no family mixing Majorana neutrinos can be distinguished from Dirac ones in the non-relativistic limit. Moreover, if CP violation is present, an effective phase arises in the effective mixing matrix but, due to a symmetry of the Majorana fields, this cannot be univocally determined

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1 Introduction

An open question of the Modern Elementary Particle Physics is the problem of neutrino mass. In the Standard Model by Glashow, Weinberg and Salam [1] neutrinos are massless but there is no fundamental reason to believe in that (on the contrary to that happens for photons, for which the electromagnetic gauge invariance forces them to be massless), and in particular the structure of the Standard Model remains unchanged if we introduce a neutrino mass term in the electroweak lagrangian. On the other hand, there are some experimental indications (even if not conclusive on this argument) [2] that neutrinos could be massive, such as the solar neutrino problem [3], the atmospheric neutrino anomaly [4] and so on. Moreover, in Grand Unified Theories it is quite natural for neutrinos to acquire a non vanishing mass [5].

In general, if neutrinos are massive, they could be of the "Dirac-type" or of the "Majorana-type". For Dirac neutrinos, in strict analogy to that happens for the other fermions present in the Standard Model, the particle states are different from the anti-particle states, so in addition to ν_L and ν_R^C, revealed by weak interactions, there must be also the states ν_R and ν_L^C, which however are sterile for weak interaction. In a lagrangian description, a Dirac neutrino of 4-momentum k_μ propagating freely in vacuum is described by

$$\mathcal{L} \;=\; \overline{\nu}_{\mathcal{L}} \; \not{k} \nu_{\mathcal{L}} + \overline{\nu}_{\mathcal{R}} \; \not{k} \nu_{\mathcal{R}} - \overline{\nu}_{\mathcal{L}} \mathcal{M}_{\mathcal{D}} \nu_{\mathcal{R}} - \overline{\nu}_{\mathcal{R}} \mathcal{M}_{\mathcal{D}}^{\dagger} \nu_{\mathcal{L}} \tag{1}$$

where $\nu^T \;=\; (\nu_e, \nu_\mu, \nu_\tau)$ are the flavour-eigenstate neutrino fields and M_D is a 3x3 Dirac mass matrix.

A more economical way for accounting neutrinos mass is that of considering neutrinos to be Majorana particles, for which the particle states coincide (up to a phase factor) with the anti-particle states, so there is no need for introducing sterile ν_R, ν_L^C states. In this case the lagrangian term (1) is substituted by

$$\mathcal{L} \;=\; \overline{\nu}_{\mathcal{L}} \; \not{k} \nu_{\mathcal{L}} + \overline{\nu}_{\mathcal{R}}^{\mathcal{C}} \; \not{k} \nu_{\mathcal{R}}^{\mathcal{C}} - \overline{\nu}_{\mathcal{L}} \mathcal{M}_{\mathcal{M}} \nu_{\mathcal{R}}^{\mathcal{C}} - \overline{\nu}_{\mathcal{R}}^{\mathcal{C}} \mathcal{M}_{\mathcal{M}}^{\dagger} \nu_{\mathcal{L}} \tag{2}$$

where now M_M is a 3x3 Majorana mass matrix that, from the fermionic anti-commutation rules of the neutrino fields, must be symmetric [8].

A more general lagrangian is that with both the Dirac mass term in (1) and the Majorana one in (2), whose diagonalization leads again to Majorana fields; however, for simplicity, in this paper we consider pure Majorana neutrinos with the lagrangian (2).

In general, if the mass matrices (Dirac or Majorana) are non diagonal (in analogy to that happens in the quark sector) lepton flavour violation phenomena can occur, between which neutrino oscillations seem to be the most likely to detect experimantally the effects of a non vanishing neutrino mass [6]. Anyhow, the properties of Dirac and Majorana neutrinos are completely different. For example, in the Majorana case the CPT invariance forces neutrinos to have zero diagonal electromagnetic moments, while this doesn't happen for Dirac particles. Further, while for Dirac neutrinos the global lepton number is a conserved quantity, for Majorana neutrinos this is no longer true, and phenomena such as Pontecorvo neutrino-antineutrino oscillations [7] and neutrinoless double beta decay

[2],[8] can arise. Moreover, if CP invariance in the leptonic sector is violated, as in the quark sector, the number of physical CP breaking phases in the mixing matrix depends on the Dirac or Majorana nature of neutrinos and, in particular, for Majorana neutrinos there is one CP breaking phase also in the case of mixing of only two families [9].

We stress that the knowledge of neutrino properties is of relevant importance not only for particle physics, but also for astrophysics and cosmology. In fact, for example, the most abundant neutrino sources are the active stars, in which neutrinos are produced by nuclear processes [10]; due to their weak interaction, these neutrinos play a fundamental role in the mechanism of star cooling. However, to this purpose, the study of neutrino interactions with the stellar matter (and in particular the study of coherent effects of matter on neutrino propagation) is necessary.

In this paper we want to extend the recent analysis [11] made on Dirac neutrino propagation in magnetized matter (such as that present in supernovæ, for example) to the case of Majorana neutrinos and to point out the relevant differences of behaviour.

The paper is organized as follows. In section 2 we study the propagation of Majorana neutrinos in matter in the case of no family mixing for pointing out the main properties. In section 3 the coherent effects on neutrino (flavour) oscillations is considered in the framework of CP invariance, while in section 4 the insorgence of "effective" phases in the case of CP violation is studied. Finally, in section 5 conclusions are outlined.

2 Dispersion relation and the free propagation condition

Let us first consider the propagation of non-mixed Majorana neutrinos in matter, so that the mass matrix is diagonal and we can consider one neutrino flavour at a time. From the lagrangian in Eq.(2) follows that in vacuum the equations of motion governing neutrino propagation are

$$(\not{k} - m)\, \nu_L + (\not{k} - m)\, \nu_R^C = 0 \tag{3}$$

When neutrinos propagate in matter, a way for taking into account their interactions with the particles in the background is to replace the mass m in Eq.(3) with the complete self-energy, so the equations of motions can be written in the form

$$(\not{k} - m + \Sigma_\nu)\, \nu_L + (\not{k} - m + \Sigma_{\overline{\nu}})\, \nu_R^C = 0 \tag{4}$$

For propagation in a medium with a magnetic field $\vec{B}$, the calculation of Σ_ν proceeds with the evaluation of the Feynman diagrams in Fig.1. At order G_F, the self-energy term Σ_ν can be cast in the form [11],[12]

$$\Sigma_\nu = b_L \not{u} + c_L \not{B} \tag{5}$$

where u_μ is the medium 4-velocity (for simplicity we consider the rest frame of the medium, in which $u_\mu = (1, \vec{0})$) and $B_\mu = \frac{1}{2}\epsilon_{\mu\nu\alpha\beta} u^\nu F^{\alpha\beta}$ (in the rest frame $B_\mu = (0, \vec{B})$. The

coefficients b_L,c_L are given by [14], [11]

$$b_L^{\nu_e} \simeq -\sqrt{2}G_F(N_e - \frac{1}{2}N_n) \tag{6}$$

$$b_L^{\nu_\mu,\nu_\tau} \simeq \frac{G_F}{\sqrt{2}}N_n \tag{7}$$

and

$$c_L^{\nu_e} \simeq 0 \tag{8}$$

$$c_L^{\nu_\mu,\nu_\tau} \simeq -\frac{eG_F}{\sqrt{2}}\frac{(3\pi^2 N_e)^{\frac{1}{3}}}{\pi^2} \tag{9}$$

for a degenerate Fermi gas, while

$$c_L^{\nu_e} \simeq -\frac{3eG_F}{4\sqrt{2}}\frac{N_e}{m_e^2} \tag{10}$$

$$c_L^{\nu_\mu,\nu_\tau} \simeq +\frac{3eG_F}{4\sqrt{2}}\frac{N_e}{m_e^2} \tag{11}$$

for a classical non-relativistic plasma. Here N_e,N_n are the number density of electrons and neutrons respectively, and we have assumed the medium to be electrically neutral.

The Feynman diagrams for $\Sigma_{\overline{\nu}}$ are showed in Fig.2; from [11],[12],[14] it is easy to deduce that $\Sigma_{\overline{\nu}}$ takes the form

$$\Sigma_{\overline{\nu}} = -\, b_L \,\not{u} \,-\, c_L \,\not{B} \tag{12}$$

where b_L,c_L are again given by Eqs.(6)-(11) [1] Then, the equation of motion can be written as

$$\left(\omega \,+\, b_L \,+\vec{\sigma}\cdot\left(\vec{k} \,+\, c_L\,\vec{B}\right)\right)\,\nu_L = m\,\nu_R^C \tag{13}$$

$$\left(\omega \,-\, b_L \,-\vec{\sigma}\cdot\left(\vec{k} \,-\, c_L\,\vec{B}\right)\right)\,\nu_R^C = m\,\nu_L \tag{14}$$

where ω and $\vec{k}$ are respectively the neutrino energy and momentum, while σ_i are the Pauli matrices. Introducing the helicity eigenstates ϕ_λ [13]

$$\frac{\vec{\sigma}\cdot\vec{k}}{|\vec{k}|}\,\phi_\lambda = \lambda\,\phi_\lambda \tag{15}$$

with $\lambda = \pm 1$, because the interaction here considered between neutrinos and the particles in the medium preserves neutrino helicity [11], we can write $\nu_L = \eta_1\,\phi_\lambda$ and $\nu_R^C = \eta_2\,\phi_\lambda$ with η_1,η_2 satisfying

$$\left(\omega \,+\, b_L \,+\lambda k \,+\lambda c_L\,\frac{\vec{k}\cdot\vec{B}}{k} \,-\, m\left(\omega \,-\, b_L \,-\lambda k \,+\lambda c_L\,\frac{\vec{k}\cdot\vec{B}}{k}\right)^{-1} m\right)\,\eta_1 \simeq 0 \tag{16}$$

[1] $\Sigma_{\overline{\nu}}$ can be obtained from Σ_ν multiplying every neutrinic quantity appearing in that by its phase acquired under charge coniugation (in particular $\gamma_\mu\,L \to -\gamma_\mu\,R$) and with the substitution $k_\mu \to -k_\mu$. Anyhow, at order G_F, the self-energy does not depend explicitely on k_μ, so that b_L,c_L remain unchanged.

$$\eta_2 \simeq \left(\omega - b_L - \lambda\left(k - c_L \frac{\vec{k}\cdot\vec{B}}{k}\right)\right)^{-1} m\,\eta_1 \tag{17}$$

From (16) we then deduce that for $\lambda = -1$ the dispersion relation reads

$$\omega^2 - k^2 = m^2 - 2k\,b_L + 2\omega\,c_L \frac{\vec{k}\cdot\vec{B}}{k} \tag{18}$$

while for $\lambda = +1$ the signs of b_L,c_L are reversed. The solution of Eq. (18) are the eigen-energies of propagating neutrino eigen-modes.

Let us now compare this result with the analogous dispersion relation found for Dirac neutrinos [11]:

$$\omega^2 - k^2 = m^2 - b_L(\omega + k) + c_L(\omega + k)\frac{\vec{k}\cdot\vec{B}}{k} \tag{19}$$

From this, it is immediately evident that the two dispersion relations become equal for ultrarelativistic neutrinos. We stress, however, that this conclusion is valid only at order G_F. To stop at this order in the Fermi coupling constant is sufficient in all physical situations in which the background is CP-asymmetric [14], such as stellar interior. As a consequence, there is no physical difference in the behaviour of Dirac and Majorana neutrinos produced in active stars and interacting with the matter inside. [2] We also observe that for non-relativistic neutrinos the obtained results do not apply to extremely low energy particles such as relic neutrinos present in our epoch. In fact, in this case the neutrino momentum $\vec{k}$ (of the order of $10^{-2} \div 10^{-4}\,eV$ for relic neutrinos [15]) is no longer conserved in the interaction with the medium, so that the same dispersion relation formalism loses of sense.

Let us note that when higher order than G_F become important, as for example in the Early Universe, where the particle-antiparticle symmetry is believed to be present, the propagation of Dirac neutrinos would be quantitatively different from that of Majorana neutrinos.

From (18) we also derive the modifications to the free propagation condition [11]

$$b_L\,k = c_L\,\omega\,\frac{\vec{k}\cdot\vec{B}}{k} \tag{20}$$

If this relation is realized (see [11] for the conditions under which this happens for ultra-relativistic neutrinos) the coherent effects dur properly to the magnetic field are opposite to those generated in the absence of the field, so that Majorana neutrinos can propagate freely in matter as in vacuum.

[2] In the framework of the Standard Model the intrinsic magnetic moment interaction of Dirac and Majorana neutrinos are different, but extremely small [8] compared to the effects here considered

3 CP conserving Majorana neutrino matter oscillations

Let us now examine the case in which the Majorana mass matrix in (2) is non-diagonal (but real), so that neutrino flavour oscillations can occur; for simplicity, we consider the mixing between only two generations (for example e,μ).

Now, the eigen-energies of $\lambda = -1$ neutrinos are the solution of the eigenvalue equation

$$det \left(\omega \,+\, b_L \,-\, k \,-\, c_L \frac{\vec{k}\cdot\vec{B}}{k} - \, m \left(\omega \,-\, b_L \,+\, k \,-\, c_L \frac{\vec{k}\cdot\vec{B}}{k} \right)^{-1} m \right) \,=\, 0 \qquad (21)$$

where b_L,c_L,m are matrices in the flavour space that, in the mass eigenstate basis, have the form [11]

$$m \,=\, \begin{pmatrix} m_1 & 0 \\ 0 & m_2 \end{pmatrix} \qquad (22)$$

$$b_L \,=\, \begin{pmatrix} b_L^W\, cos^2\theta \,+\, b_L^Z & -\, b_L^W\, sin\,\theta\, cos\,\theta \\ -\, b_L^W\, sin\,\theta\, cos\,\theta & b_L^W\, sin^2\theta \,+\, b_L^Z \end{pmatrix} \qquad (23)$$

$$c_L \,=\, \begin{pmatrix} c_L^W\, cos^2\theta \,+\, c_L^Z & -\, c_L^W\, sin\,\theta\, cos\,\theta \\ -\, c_L^W\, sin\,\theta\, cos\,\theta & c_L^W\, sin^2\theta \,+\, c_L^Z \end{pmatrix} \qquad (24)$$

Here, θ is the vacuum mixing angle while b_L^W,c_L^W and b_L^Z,c_L^Z are respectively the charged-current and neutral-current contribution to the coefficients in the self-energy. After simple algebra, in the approximation of non-exceptionally dense medium, we arrive at the following eigenvalue equation [3]

$$\left(\omega \,-\, \overline{k_1} \,-\, \frac{m_1^2}{\omega \,+\, k_1} \right) \left(\omega \,-\, \overline{k_2} \,-\, \frac{m_2^2}{\omega \,+\, k_2} \right) \,+$$

$$+\; (\omega \,-\, \overline{k_1})(\omega \,-\, \overline{k_2}) \left(\left(1 \,-\, \frac{S_{12}^2}{(\omega \,+\, k_1)(\omega \,+\, k_2)} \right) \left(1 \,-\, \frac{D_{12}^2}{(\omega \,-\, \overline{k_1})(\omega \,-\, \overline{k_2})} \right) \,-\, 1 \right) \,+$$

$$+\; \frac{2\, m_1\, m_2\, S_{12}\, D_{12}}{(\omega \,+\, k_1)(\omega \,+\, k_2)} \,=\, 0 \qquad (25)$$

where the matrices S, D are given respectively by $b_L + c_L\, \frac{\vec{k}\cdot\vec{B}}{k}$, $b_L - c_L\, \frac{\vec{k}\cdot\vec{B}}{k}$ and $k_1 \,=\, k - S_{11}$, $k_2 \,=\, k - S_{22}$, $\overline{k_1} \,=\, k - D_{11}$, $\overline{k_2} \,=\, k - D_{22}$. Compare this result with that analogous for Dirac neutrinos [11]

$$\left(\omega \,-\, \overline{k_1} \,-\, \frac{m_1^2}{\omega \,+\, k} \right) \left(\omega \,-\, \overline{k_2} \,-\, \frac{m_2^2}{\omega \,+\, k} \right) \,-\, D_{12}^2 \,=\, 0 \qquad (26)$$

[3] Neglecting damping effects, the energy eigenvalues completely determine neutrino flavor oscillations. In fact, all the physically interesting quantities (survival probability, resonance and no-resonance conditions and so on) can be expressed in terms of these quantities, as can be seen, for example, in [11]

Let us observe that in astrophysical environments the terms

$$\frac{b_L^W \, - \, c_L^W \, \frac{\vec{k}\cdot\vec{B}}{k}}{\omega \, + k} \qquad , \qquad \frac{b_L^Z \, - \, c_L^Z \, \frac{\vec{k}\cdot\vec{B}}{k}}{\omega \, + k} \tag{27}$$

are completely negligible, so that we arrive at the conclusion that Majorana neutrino matter oscillations are indistinguishable from Dirac ones, irrispective of the relativistic properties of the neutrinos themselves. Obviously, as in section 2, this conclusion is not applyable to oscillations in the Early Universe.

4 Effective CP violating phases

Finally, let us consider the general case of a complex non-diagonal Majorana mass matrix arising when CP invariance in the leptonic sector is violated. Now, the mixing matrix V is unitary and, in general, for two generations three phases are present, two of which physical meaningless because they can be absorbed [8] by the charged lepton fields in the weak charged interaction lagrangian term. Then the mixing matrix V can be cast in the form [9]

$$V \;=\; \begin{pmatrix} cos\,\theta & -\,sin\,\theta\,e^{i\delta} \\ sin\,\theta\,e^{-i\delta} & cos\,\theta \end{pmatrix} \tag{28}$$

Considering matter oscillations, the eigenvalue equation is again of the form (21), but now the matrix b_L, in the mass eigenstate basis, is given by

$$b_L \;=\; \begin{pmatrix} b_L^W\,cos^2\theta \,+\, b_L^Z & -\,b_L^W\,sin\,\theta\,cos\,\theta\,e^{i\delta} \\ -\,b_L^W\,sin\,\theta\,cos\,\theta\,e^{-i\delta} & b_L^W\,sin^2\theta \,+\, b_L^Z \end{pmatrix} \tag{29}$$

and analogously for c_L. For definiteness, let us consider only ultrarelativistic neutrinos. The eigenvalue problem (21) is equivalent to the diagonalization of the hamiltonian [13]

$$H \;\simeq\; k \,+\, \frac{m^2}{2k} \,-\, b_L \,+\, c_L\,\frac{\vec{k}\cdot\vec{B}}{k} \tag{30}$$

that, disregarding terms proportional to the identity matrix which play no role in neutrino oscillations, in the flavour eigenstate basis takes the form

$$H_{flav} \;=\; \frac{\Delta m^2}{2k} \begin{pmatrix} sin^2\theta & -sin\,\theta\,cos\,\theta\,e^{i\delta} \\ -\,sin\,\theta\,cos\,\theta\,e^{-i\delta} & cos^2\theta \end{pmatrix} \,-\, \left(b_L^W \,-\, c_L^W\,\frac{\vec{k}\cdot\vec{B}}{k} \right) \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix} \tag{31}$$

This is an hermitian matrix that can be diagonalized by a unitary effective mixing matrix $\tilde{V_m}$ ($\tilde{V_m}^\dagger\,H_{flav}\,\tilde{V_m} \;=\; H_m$ with H_m diagonal and real) whose general form is [9]

$$\tilde{V_m} \;=\; P\,V_m \;=\; \begin{pmatrix} e^{i\theta_1} & 0 \\ 0 & e^{i\theta_2} \end{pmatrix} \begin{pmatrix} cos\,\theta_m & -\,sin\,\theta_m\,e^{i\alpha} \\ sin\,\theta_m\,e^{-i\alpha} & cos\,\theta_m \end{pmatrix} \tag{32}$$

Obviously, the phases θ_1,θ_2 don't produce separately observable effects, but only their difference $\beta = \theta_2 - \theta_1$ could. In fact, $P^\dagger H_{flav} P$ has the same form as H_{flav} but with δ replaced by $\delta + \beta$. Furthermore, regarding neutrino oscillations, only one effective phase δ_m plays a physical role. In fact, the diagonalization equation which determine the effective mixing angle θ_m and the effective phase δ_m are

$$\left(\frac{\Delta m^2}{2k} \cos 2\theta + b_L^W - c_L^W \frac{\vec{k}\cdot\vec{B}}{k}\right) \sin 2\theta_m =$$

$$= \frac{\Delta m^2}{2k} \sin 2\theta \left(\cos^2\theta_m \, e^{i(\delta+\delta_m)} - \sin^2\theta_m \, e^{-i(\delta+\delta_m)}\right) \tag{33}$$

$$\cos^2\theta_m \, e^{i(\delta+\delta_m)} - \sin^2\theta_m \, e^{-i(\delta+\delta_m)} = \cos^2\theta_m \, e^{-i(\delta+\delta_m)} - \sin^2\theta_m \, e^{i(\delta+\delta_m)} \tag{34}$$

with $\delta_m = \beta - \alpha$. From these we get that the effective mixing angle θ_m is unaffected respect to the case of CP conservation, and this follows from the fact that the effective phase is exactly opposite to the intrinsic (vacuum) one:

$$tg\, 2\theta_m = \frac{\frac{\Delta m^2}{2k} \sin 2\theta}{\frac{\Delta m^2}{2k} \cos 2\theta + b_L^W - c_L^W \frac{\vec{k}\cdot\vec{B}}{k}} \tag{35}$$

$$\delta_m = -\delta \tag{36}$$

We now give a physical interpretation of the obtained results.

In vacuum, there is no way to observe CP violating phases in flavour oscillations, because a CP violation means that there is a difference in the physical behaviour of neutrinos and antineutrinos, so that if we observe only transistions between neutrinos (even if of different flavour) this difference can not be detectable. The same is true for flavour matter oscillations (even if neutrinos and antineutrinos have different refractive index), so the effective mixing angle θ_m must have no changes and this explains (35).

Note, however, that (36) is a bit more general condition, and it is surprising the fact that the effective CP violating phase δ_m determined by (36) does not determine univocally the effective phase α present in the mixing matrix (32). This result is a consequence of the fact that the physical meaningful effective mixing matrix has the same form as the vacuum mixing matrix, i.e. it is V_m, not $\tilde{V_m}$, or, in other words, that the matrix P has no physical effect. In fact, as already noted, it has no effect on the hamiltonian ($\delta \to \delta + \beta$), but this is not sufficient because it must have no effect also on the Majorana matter mass eigenstate neutrino fields. That happens because these fields can be written in the form

$$\begin{pmatrix} \nu_{1m} \\ \nu_{2m} \end{pmatrix} = V_m^\dagger P^\dagger \begin{pmatrix} \nu_{eL} \\ \nu_{\mu L} \end{pmatrix} + V_m^T P \begin{pmatrix} \nu_{eR}^C \\ \nu_{\mu R}^C \end{pmatrix} \tag{37}$$

and the phases present in P can be absorbed *simultaneously* by ν_L and ν_R^C, redefining the physical states as

$$\left(\begin{array}{c} \nu_{eL} \\ \nu_{\mu L} \end{array} \right) \rightarrow P^\dagger \left(\begin{array}{c} \nu_{eL} \\ \nu_{\mu L} \end{array} \right) = \left(\begin{array}{c} e^{-i\theta_1} \, \nu_{eL} \\ e^{-i\theta_2} \, \nu_{\mu L} \end{array} \right) \tag{38}$$

$$\left(\begin{array}{c} \nu_{eR}^C \\ \nu_{\mu R}^C \end{array} \right) \rightarrow P \left(\begin{array}{c} \nu_{eR}^C \\ \nu_{\mu R}^C \end{array} \right) = \left(\begin{array}{c} e^{i\theta_1} \, \nu_{eR}^C \\ e^{i\theta_2} \, \nu_{\mu R}^C \end{array} \right) . \tag{39}$$

5 Conclusions

In this paper we have analyzed the behaviour of pure Majorana neutrinos when they pass through a dense medium (typically a stellar medium) and outlined the main differences with the behaviour of Dirac neutrinos.

The self-energy in matter of neutrino and antineutrino states are different (compare (5) with (12)). However, this brings to observable differences in the propagation of Majorana neutrinos in stellar matter with respect to the Dirac case only for non-relativistic neutrinos (at first order in the Fermi coupling constant). It is anyhow remarkable the fact that Majorana neutrino matter oscillations are completely indistinguishable from Dirac ones in astrophysical environments (both for ultrarelativistic anf for non-relativistic neutrinos), due to the smallness of the effective potential (27).

Finally, we have considered the case of a possible CP (intrinsic) violation that can be present, for Majorana neutrinos, also in the mixing of only two generations. Obviously this has no effect on neutrino flavour conversion (both in vacuum and in matter) but should have on neutrino-antineutrino oscillations. We have then analyzed the properties of the effective unitary mixing matrix and have shown that the effective phase α present in it, due to the invariance for phase change of the Majorana matter eigenstate neutrino fields (37), is not univocally determined.

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